

# Extended Zee model for Neutrino Mass, Leptogenesis and Sterile Neutrino like Dark Matter

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We propose an extension of the standard model with a  $U(1)_{B-L}$  global symmetry that accommodates radiative neutrino masses along with dark matter and leptogenesis. The observed matter antimatter asymmetry of the universe is generated through the leptogenesis route keeping the  $U(1)_{B-L}$  symmetry intact. The  $B-L$  global symmetry is then softly broken, providing the sub-eV neutrino masses. The model then incorporates a MeV scale sterile neutrino like dark matter.

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## I. INTRODUCTION

In recent times, there have been growing interests on models of neutrino mass, dark matter and baryon asymmetry [1, 2]. While the low energy neutrino oscillation data [3, 4, 5] indicate that at least two of the physical left-handed (LH) neutrinos have tiny masses and therefore they mix among themselves, the galactic rotation curve [6], gravitational lensing [7] and large scale structure [8] strongly demand that there exist non-baryonic dark matter in the present universe, which has only matter (visible plus dark) components. The antimatter components in the present universe are vanishingly small. This implies that the asymmetry between matter and antimatter components is maximal today. Currently this asymmetry has been precisely measured by the Wilkinson Microwave Anisotropy Probe (WMAP) [9] and is given by

$$\left( \frac{n_B - n_{\bar{B}}}{n_\gamma} \right)_0 \equiv \left( \frac{n_B}{n_\gamma} \right)_0 = 6.1_{-0.2}^{+0.3} \times 10^{-10}, \quad (1)$$

where  $n_B$  is the baryon density and  $n_\gamma$  is the primordial photon density.

Despite the success of the standard model (SM), it can not explain any of the above phenomena: non-zero neutrino mass, existence of dark matter and matter-antimatter asymmetry of the universe. Explanation of all of these phenomena requires an extension of SM. The observed neutrino oscillation experiments deal only with the mass square differences of the neutrinos, and hence, they leave an ambiguity on the nature of neutrinos to be either Dirac or Majorana. If the neutrinos are assumed to be Majorana particles, the sub-eV neutrino masses can naturally be generated through the celebrated see-saw mechanisms: type-I [10] and type-II [11]. In either case, the neutrino mass is suppressed by the scale of lepton (L) number violation. The smallness of the neutrino mass could also be explained naturally without invoking any large lepton number violating scale by the radiative mechanisms [12]. The loop factor then introduces the suppression required to keep

the neutrino masses small. However, these models have the generic problem of explaining the baryon asymmetry of the universe, since the fast lepton number violation required to generate the neutrino masses, also erases any baryon asymmetry of the universe [13]. In this note we aim to solve this problem and generate the neutrino masses and leptogenesis simultaneously in a radiative model. The lepton asymmetry is generated without any B-L violation, similar to leptogenesis in models of Dirac neutrinos [14, 15]. The present model also accommodates a sterile neutrino like dark matter candidate. In contrast to the generic radiative neutrino mass models, where the masses of the fields propagating in the loop are expected to be at the TeV scale, in the present model the masses of the fields propagating in the loop could be as high as  $10^{10}$  GeV or so. As a result a leptogenesis could be possible from the out-of-equilibrium decay of these heavy charged scalars in the early universe.

The paper is arranged as follows. In section II we propose a radiative neutrino mass model which has an exact B-L symmetry to begin with. The B-L symmetry is then softly broken to give rise neutrino masses. Section III is devoted to estimate neutrino masses originating from the soft B-L violation. In section IV we calculate lepton asymmetry from a conserved B-L symmetry. In section V we estimate the lifetime of a decaying sterile neutrino like dark matter. Finally section VI concludes.

## II. EXTENDED ZEE MODEL WITH CONSERVED B-L SYMMETRY

Within the SM neutrinos are massless. This is because  $B-L$  is an exact symmetry of SM. So any radiative mechanism within SM can not give rise to neutrino masses. One of the simplest extensions of SM is Zee model where SM is extended with a charged scalar, say  $\eta^-$ , and a second doublet Higgs, say  $\phi_b$ , to generate neutrino masses through one loop radiative correction. The antisymmetric couplings of  $\eta^-$  to Higgses ( $\mu_{ab}\eta^-\phi_a\phi_b$ ) and leptons ( $f_{ij}\eta^-\ell_i\ell_j$ ) then together violate lepton number by two units and neutrinos acquire masses through one loop radiative correction. Since the couplings of the Higgses and the charged leptons to the charged scalar

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is antisymmetric, the diagonal elements of the neutrino mass matrix vanishes identically in a flavor basis. The lepton number violating couplings that appear in the neutrino masses, also give rise to lepton number violating interactions in the early universe. The observed neutrino masses would then imply that the lepton number violating interactions in the early universe is too fast, which will wash out [13] any matter–antimatter asymmetry of the universe in the presence of the sphalerons before the electroweak phase transition.

We propose an extension of the Zee model, which can simultaneously explain the observed neutrino masses, the baryon asymmetry of the universe and also provide a candidate for dark matter of the universe. We extend the Zee model with an additional charged scalar  $\chi^-$  and three sterile neutrino like fields  $N_L$ , and impose a  $U(1)_{B-L}$  global symmetry. The particle content of the model and their quantum numbers are presented in table I. The relevant terms in the Lagrangian are

TABLE I: Particle content and their quantum numbers.

Particle Content	$SU(3)_C \times SU(2)_L \times U(1)_Y$	$U(1)_{B-L}$
$\ell_L$	(1,2,-1)	-1
$e_R^-$	(1,1,-2)	-1
$\phi_a, \phi_b$	(1,2,1)	0
$\chi^-$	(1,1,-2)	-2
$\eta_a^-$	(1,1,-2)	0
$N_L$	(1,1,0)	-1

then given as:

$$\mathcal{L} \supseteq M_\chi^2 \chi^\dagger \chi + M_\eta^2 \eta^\dagger \eta + f_{ij} \chi^\dagger \ell_{iL} \ell_{jL} + (\mu)_{ab} \eta \phi_a \phi_b + h_{ij} \eta^\dagger \overline{N_{iL}} e_{jR} + Y_{ii}^a \overline{\ell_{iL}} \phi_a e_{iR} + h.c. \quad (2)$$

where  $i, j = e, \mu, \tau$  are family indices and  $a, b$  corresponds to two Higgs doublets. Since  $N_L$  is electrically neutral it can have Majorana masses  $M_{N_L}$ . In Eq. (2) we have assumed that  $\phi_a$  couples only to leptons and  $\phi_b$  couples only to quarks [16] apart from their self-interactions. However, this can be achieved by using an additional  $Z_2$  symmetry. We also assume that the couplings of leptons to  $\phi_a$  is diagonal similar to ref. [16]. As a result there is no tree level flavor changing processes induced by  $Y_{ij}$  in the lepton sector.

As the fields  $\chi^-$  and  $\eta^-$  are charged they can not acquire any vacuum expectation value (VEV). The scalar potential involving the  $\phi_a$  and  $\phi_b$  then can be given as:

$$V(\phi_a, \phi_b) = -M_a^2 |\phi_a|^2 - M_b^2 |\phi_b|^2 + M_{ab}^2 \phi_a^\dagger \phi_b + \lambda_a |\phi_a|^4 + \lambda_b |\phi_b|^4 + \lambda_{ab} |\phi_a|^2 |\phi_b|^2 + \lambda'_{ab} |\phi_a^\dagger \phi_b|^2 + \frac{\lambda''_{ab}}{2} [(\phi_a^\dagger \phi_b)^2 + h.c.] \quad (3)$$

where  $M_a^2, M_b^2 > 0$ . The stability of the potential requires  $\lambda_a, \lambda_b > 0$  and  $\lambda_{ab} > -\sqrt{\lambda_a \lambda_b}$ . The VEV of Higgses  $\phi_a$  and

$\phi_b$  can be given as

$$\langle \phi_a \rangle = v_a \quad \text{and} \quad \langle \phi_b \rangle = v_b \quad (4)$$

Then the two VEVs are related by

$$v = \sqrt{v_a^2 + v_b^2} = 174 \text{ GeV}, \quad v_a = v \sin \beta \quad \text{and} \quad v_b = v \cos \beta, \quad (5)$$

where  $\beta = \tan^{-1}(v_a/v_b)$ .

### III. SOFT B-L VIOLATION AND RADIATIVE NEUTRINO MASSES

The  $U(1)_{B-L}$  symmetry is allowed to be broken softly by

$$\mathcal{L}_{\text{soft}} = m^2 \eta^\dagger \chi + h.c. \quad (6)$$

As a result there is a mixing between  $\eta^-$  and  $\chi^-$ . In the flavor basis the mass matrix of  $\eta^-$  and  $\chi^-$  is then given by

$$\mathcal{M}^2 = \begin{pmatrix} M_\eta^2 & m^2 \\ m^2 & M_\chi^2 \end{pmatrix}. \quad (7)$$

Diagonalising the above mass matrix we get the eigenvalues:

$$M_{\eta', \chi'}^2 = \frac{(M_\eta^2 + M_\chi^2) \pm \sqrt{(M_\eta^2 - M_\chi^2)^2 + 4m^4}}{2} \quad (8)$$

corresponding to the mass eigenstates  $\eta' = \cos \theta \eta^- + \sin \theta \chi^-$  and  $\chi' = \cos \theta \chi^- - \sin \theta \eta^-$ , where

$$\theta = \frac{1}{2} \tan^{-1} \left( \frac{2m^2}{M_\chi^2 - M_\eta^2} \right). \quad (9)$$

Through the mixing between  $\eta^-$  and  $\chi^-$  lepton number is violated by two units. As a result neutrino mass is generated through the one loop radiative correction diagram as shown in fig. 1. From fig. 1, the neutrino mass can be estimated to be

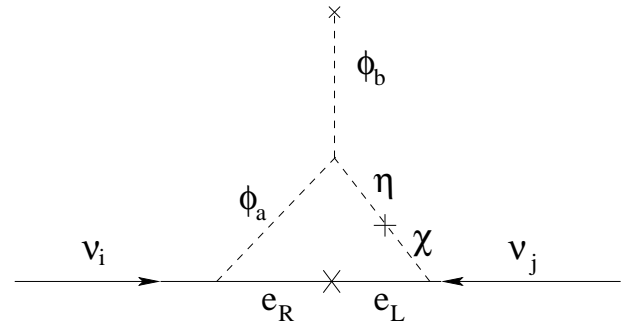


FIG. 1: One loop radiative correction diagram for neutrino masses arises through the mixing between  $\eta$  and  $\chi$ .

$$(M_\nu)_{ij} = (Y_{ii}^a M_i - Y_{jj}^a M_j) f_{ij} (m^2 \mu_{ab} v_b) \cdot |I| \quad (10)$$

where  $M_i$  and  $M_j$  are the diagonal charged lepton mass matrix and the integral:

$$I = \int \frac{d^4 q}{(2\pi)^4} \frac{1}{q^2 - M_\chi^2} \frac{1}{q^2 - M_\eta^2} \frac{1}{q^2 - M_a^2} \frac{1}{q^2} \\ = \frac{i}{16\pi^2} \left[ \frac{M_a^2 \ln(M_\chi^2/M_\eta^2) + M_\chi^2 \ln(M_\eta^2/M_a^2) + M_\eta^2 \ln(M_a^2/M_\chi^2)}{(M_\chi^2 - M_\eta^2)(M_\chi^2 - M_a^2)(M_\eta^2 - M_a^2)} \right] \quad (11)$$

From Eq. (10) it is clear that  $M_\nu$  is symmetric with respect to the family indices  $i$  and  $j$  and the diagonal elements of  $M_\nu$  are also zero. This is because of our assumption about the coupling of  $\phi_a$  to the leptons. However, in general, this is not true.

If we assume that  $M_\eta, M_\chi \gg M_a, M_b$ , then the neutrino mass is simply

$$(M_\nu)_{ij} \simeq \frac{1}{16\pi^2} (M_i^2 - M_j^2) f_{ij} \frac{\mu_{ab} m^2 \cot \beta}{(M_\chi^2 - M_\eta^2)} \\ \times \left[ \frac{\ln(M_\eta^2/M_a^2)}{M_\eta^2} - \frac{\ln(M_\chi^2/M_b^2)}{M_\chi^2} \right] \\ = M_0 F_{ij}, \quad (12)$$

where

$$M_0 = \frac{1}{16\pi^2} \frac{M_\tau^2 \mu_{ab} m^2}{(M_\chi^2 - M_\eta^2)} \cot \beta \left[ \frac{\ln(M_\eta^2/M_a^2)}{M_\eta^2} - \frac{\ln(M_\chi^2/M_b^2)}{M_\chi^2} \right] \quad (13)$$

and  $F_{ij} = f_{ij} \delta_{ij}$ , where

$$\delta_{ij} = \frac{M_i^2 - M_j^2}{M_\tau^2}. \quad (14)$$

We now estimate the magnitude of front factor  $M_0$  by using the following sample set:

$$M_\eta = 3 \times 10^{10} \text{ GeV}, \quad M_\chi = 5 \times 10^{10} \text{ GeV} \quad \text{and} \quad M_a = 500 \text{ GeV} \quad (15)$$

We also choose the mass dimension coupling to be  $\mu = 10^{15}$  GeV and the soft B-L violation scale  $m = 10^9$  GeV. Then we get the symmetric neutrino mass matrix

$$(M_\nu)_{ij} = 0.3 \text{ eV} \cot \beta f_{ij} \delta_{ij}. \quad (16)$$

The coupling constants  $F_{ij} = f_{ij} \delta_{ij}$  can be appropriately chosen to explain the current neutrino oscillation data. Similar to the other neutrino mass models the coupling constants  $F_{ij}$  are constrained by the flavor changing processes like  $\mu \rightarrow e + \gamma$ ,  $\tau \rightarrow \mu + \gamma$ , etc....

Note that in contrast to the original Zee model here the charged scalars flowing through the loop are super heavy. Therefore, they remain in out-of-thermal equilibrium for a while in the early universe, and hence, can generate a net lepton asymmetry consistently as discussed below.

#### IV. LEPTOGENESIS FROM A CONSERVED $B - L$ SYMMETRY

As the universe expands the temperature of the thermal bath falls. As a result  $\eta^-$  will go out-of-thermal equilibrium below its mass scale. Note that  $\eta^-$  has gauge interaction:  $\eta^- \eta^+ \rightarrow B_\mu B^\mu$ , where  $B_\mu$  is the  $U(1)_Y$  gauge field. However, for  $M_\eta \gtrsim 10^{10}$  GeV the gauge interactions will remain in out-of-thermal equilibrium for a decoupled temperature, say,  $T_D \simeq M_\eta/10$ . The partial decay width of  $\eta^- \rightarrow \bar{N}_{iL} e_{jR}^-$  can be given as

$$\Gamma_\eta = \frac{1}{8\pi} |h_{ij}|^2 M_\eta, \quad (17)$$

where the family index  $i, j = e, \mu, \tau$ . At a cosmic temperature  $T \simeq M_\eta$ , if  $\Gamma_\eta$  fails to compete with the Hubble expansion parameter

$$H = 1.67 g_*^{1/2} \frac{T^2}{M_{\text{pl}}}, \quad (18)$$

where  $g_* = 106$  is the number of relativistic degrees of freedom, then  $\eta$  goes out-of-thermal equilibrium. From equations (17) and (18) then we find that the out-of-equilibrium decay of  $\eta^-$  occurs for

$$M_\eta \gtrsim 2.77 \times 10^{10} \text{ GeV} \left( \frac{h_{ij}}{10^{-3}} \right)^2. \quad (19)$$

Note that decay of  $\eta^-$  produce a pair of lepton and anti-

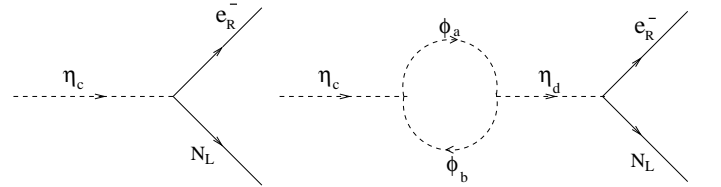


FIG. 2: Tree level and one loop self energy correction diagrams whose interference generates a net CP asymmetry.

lepton. Therefore, the decay of  $\eta$  does not produce any lepton asymmetry. However, if there are at least two  $\eta^-$  fields then there can be CP violation in the decay if  $\eta$  fields. In the presence of their interactions, the diagonal mass  $M_\eta^2$  in equation (2) is replaced by

$$\eta_c^\dagger (\mathcal{M}_+^2)_{cd} \eta_d + (\eta_c^*)^\dagger (\mathcal{M}_-^2)_{cd} \eta_d^* \quad (20)$$

where

$$\mathcal{M}_\pm^2 = \begin{pmatrix} M_{\eta_1}^2 - i\mathcal{G}_{11} & -i\mathcal{G}_{12}^\pm \\ -i\mathcal{G}_{21}^\pm & M_{\eta_2}^2 - i\mathcal{G}_{22} \end{pmatrix}, \quad (21)$$

where  $\mathcal{G}_{cd}^+ = \Gamma_{\eta_{cd}} M_{\eta_d}$ ,  $\mathcal{G}_{cd}^- = \Gamma_{\eta_{cd}}^* M_{\eta_d}$ , and  $\mathcal{G}_{cc} = \Gamma_{\eta_{cc}} M_{\eta_c}$  with  $\Gamma_{\eta_{cc}} \equiv \Gamma_{\eta_c}$ . Similarly, the interaction term  $\mu_{ab} \eta \phi_a \phi_b$  in

equation (2) should be replaced by  $(\mu_{ab})_c \eta_c \phi_a \phi_b \equiv \mu_c \eta_c \phi_a \phi_b$  and  $(\mu_{ab})_d \eta_d \phi_a \phi_b \equiv \mu_d \eta_d \phi_a \phi_b$ , where we have suppressed the symbols “ab” in  $\mu$ . Now the absorptive part of one loop self energy diagram for  $\eta_c \rightarrow \eta_d$  can be given by

$$\Gamma_{\eta_{cd}} M_{\eta_d} = \frac{1}{8\pi} \left( \mu_c \mu_d^* + M_{\eta_c} M_{\eta_d} \sum_{i,j} h_{cij} h_{dij}^* \right). \quad (22)$$

Diagonalising the mass matrix (21) one will get two mass eigenvalues corresponding to the two mass eigenstates  $\psi_1^\pm$  and  $\psi_2^\pm$ . Note that the mass eigenstates  $\psi_1^+$  and  $\psi_1^-$  (similarly  $\psi_2^+$  and  $\psi_2^-$ ) are not CP conjugate of each other even though they are degenerate mass eigenstates, while  $\eta_1^+$  and  $\eta_1^-$  (similarly  $\eta_2^+$  and  $\eta_2^-$ ) are CP conjugates of each other. Therefore, the decay of the lightest  $\psi^\pm$ , say  $\psi_1^\pm$  can generate a net CP asymmetry through the interference of tree level and self energy correction diagram as shown in fig. (2). The CP asymmetry is then given by [17],

$$\varepsilon_1 = \frac{\text{Im} \left[ (\mu_1 \mu_2^*) \sum_{ij} h_{1ij} h_{2ij}^* \right]}{16\pi^2 (M_{\eta_2}^2 - M_{\eta_1}^2)} \left[ \frac{M_{\eta_1}}{\Gamma_{\eta_1}} \right]. \quad (23)$$

Note that there is no radiative correction diagram. The decay of  $\psi_1^\pm$  does not violate L-number, since the decay of  $\eta^\pm$  does not violate lepton number. Therefore, the out-of-equilibrium-decay of  $\psi_1^\pm$  does not produce any L-asymmetry. However, the decay of  $\psi_1^\pm$ , below its mass scale, generates an equal B-L asymmetry between  $N_L$  and  $e_R$  due to the CP violation [1]. The B-L asymmetry stored in  $e_R$  is then transferred to  $e_L$  through the  $t$ -channel process  $e_R e_R^c \leftrightarrow \phi_a^0 \leftrightarrow e_L e_L^c$  as shown in the figure (3). These interactions will remain in thermal

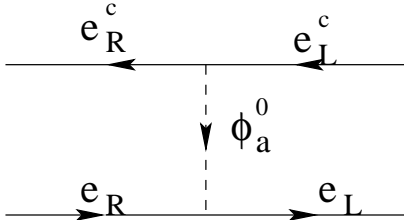


FIG. 3: The L-number conserving process which transfer the B-L asymmetry from right handed sector to the left-handed sector.

equilibrium for all three generations of charged leptons below  $10^5$  GeV and hence there will be an equal amount of  $e_R$  and  $e_L$  asymmetry. The B-L asymmetry in  $e_L$  will be converted to the baryon asymmetry of the universe before the electroweak phase transition when the sphaleron processes are in thermal equilibrium, while an equal asymmetry will remain in  $N_L$ . The two asymmetries will equilibrate with each other after the electroweak phase transition when the sphaleron processes go out-of-thermal equilibrium.

The final baryon asymmetry thus generated can be given as

$$\frac{n_B}{s} \simeq \left( \frac{28}{79} \right) \frac{\varepsilon_1}{g_* K (\ln K)^{0.6}}, \quad (24)$$

where  $K \equiv \Gamma_1/H$ ,  $\Gamma_1$  being the decay rate of  $\psi_1$ , measures the deviation from equilibrium. If  $K \ll 1$  then the baryon asymmetry is simply  $(n_B/s) \sim \varepsilon_1/g_*$ . On the other hand, if  $K > 1$ , then the final baryon asymmetry suffers from a suppression of  $1/K(\ln K)^{0.6}$ . If we assume that the CP violation is maximal, then by substituting  $M_{\eta_1}^2/(M_{\eta_2}^2 - M_{\eta_1}^2) \simeq O(1)$  and  $\mu_2/\mu_1 \simeq O(1)$  in equation (23) we get a CP asymmetry  $\varepsilon_1 \sim 10^{-6}$  for  $h_{1ij} \simeq h_{2ij} \sim 10^{-3}$ . Smaller values of  $\varepsilon_1$  can be conspired if we assume non-maximal CP violation. If we further assume that  $K = 1.0 \times 10^3$ , then the suppression factor will be  $3 \times 10^{-4}$ . As a result we will get a net baryon asymmetry  $n_B/s \sim 10^{-10}$ . The observed baryon asymmetry is then given by  $\eta = n_B/n_\gamma = 7(n_B/s)$ . An exact value of the suppression factor can be obtained by solving the required Boltzmann equation numerically which is beyond the scope of this paper.

### A. Washout Constraints

Below the mass scale of lightest  $\eta^-$ , the interaction  $h_{ij}^* \eta^- \bar{e}_{jR} N_{iL}$  is already gone out-of-thermal equilibrium. Therefore, there is no direct transfer of B-L asymmetry stored in  $N_L$  to  $e_R$ . However, in a thermal bath the scattering:  $\bar{N}_{iL} e_{jR} \rightarrow \nu_{iL} e_{jL}$  can occur through the mixing between  $\eta^-$  and  $\chi^-$  as shown in fig. (4). Note that this process violate lepton number by two units and therefore it is suppressed by the large mass scale of  $\eta^-$  and  $\chi^-$ .

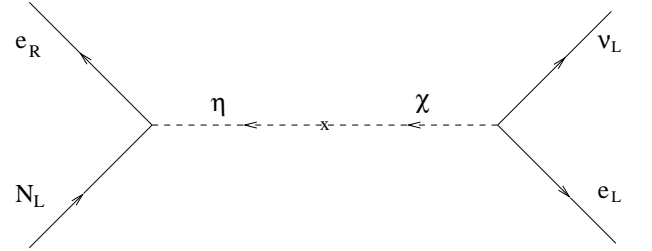


FIG. 4: Scattering of  $N_L$  and  $e_R$  through the mixing between  $\eta^-$  and  $\chi^-$ .

To ensure that, we compute the scattering cross-section at a temperature below the mass scale of lightest  $\eta^-$ :

$$\langle \sigma | v | \rangle = \frac{1}{2\pi} \frac{|h_{ij}|^2 |f_{ij}|^2 m^4 E^2}{M_\eta^4 M_\chi^4}, \quad (25)$$

where we have assumed that  $E_e = E_N = E$ . Thus from the above equation we see that below the mass scale of lightest  $\eta^-$  and  $\chi^-$ , the cross-section goes as:  $\langle \sigma | v | \rangle \propto m^4/M^8$ , assuming that  $M_\eta \approx M_\chi = M$ . The corresponding scattering rate for  $N_{iL}$ ,  $i = e, \mu, \tau$  can be given as:

$$\Gamma_{N_L} = n_{N_L}^{eq} \langle \sigma | v | \rangle \quad (26)$$

where  $n_{N_L}^{eq}$  is the equilibrium number density of  $N_L$  and it is given by

$$n_{N_L}^{eq} = \frac{2T^3}{\pi^2} \quad (27)$$

The out-of-equilibrium of the scattering process:  $\overline{N}_{iL}e_{jR} \rightarrow \nu_{iL}e_{jL}$  then requires

$$\Gamma_{N_L} \lesssim H(T \simeq M), \quad (28)$$

where  $H$  is the Hubble expansion parameter. From equation (28) we get a constraint on the soft B-L violation scale to be

$$m \lesssim 8.54 \times 10^9 \text{ GeV} \left( \frac{M}{10^{10} \text{ GeV}} \right)^5 \left( \frac{O(1)}{|f_{ij}|^2} \right) \left( \frac{10^{-6}}{|h_{ij}|^2} \right), \quad (29)$$

where we have used  $g_* = 106$ . Thus from the above equation we get a constraint:  $m < O(10^{10})$  GeV in order that the scattering will remain in out-of-thermal equilibrium above the electroweak phase transition.

### B. Low reheat temperature and Viability of leptogenesis

It is believed that the universe has gone through a period of inflation and then reheated to a uniform temperature  $T_{\text{reh}}$ . If the corresponding theory of matter is supersymmetric then  $T_{\text{reh}}$  is highly constrained by the success of Big-Bang nucleosynthesis, which could be spoiled by the overproduction of gravitino during the radiation dominated epoch [18]. For a gravitino mass of  $O(100 \text{ GeV} - 1 \text{ TeV})$ , a conservative upper bound on  $T_{\text{reh}}$  reads  $10^{6-9} \text{ GeV}$  [19]. If  $T_{\text{reh}}$  is the maximum temperature during reheating then it is difficult to create sufficiently high number densities of GUT gauge and Higgs bosons including the super heavy  $\eta^\pm$  and  $\chi^\pm$ . However, as discussed in refs. [20, 21], after the inflationary era the temperature does not rise instantaneously to  $T_{\text{reh}}$ , but rises initially to a maximum temperature  $T_{\text{max}}$  and then falls to  $T_{\text{reh}}$ . It is argued that  $T_{\text{max}}$  can be as high as  $10^3 T_{\text{reh}}$  [21]. As a result the super-heavy charged particles  $\eta^\pm$  and  $\chi^\pm$  can be easily produced through the gauge interactions:  $B_\mu B^\mu \rightarrow \eta^+ \eta^-$  and  $B_\mu B^\mu \rightarrow \chi^+ \chi^-$ . Even though  $T_{\text{max}}$  is quite higher than  $T_{\text{reh}}$ , the gravitinos are mostly produced at  $T_{\text{reh}}$  [22] and therefore,  $T_{\text{max}} \gg T_{\text{reh}}$  is no more harmful for the success of Big-Bang nucleosynthesis. Subsequently the CP violating decay of these charged particles can produce lepton asymmetry consistently as discussed above.

### V. DECAYING DARK MATTER

As  $N_1$  is neutral, it can be a candidate of dark matter. Since  $B - L$  symmetry is already broken, it can not be a stable dark

matter. It will decay through the three body process:  $N_1 \rightarrow e^- e^+ \nu$ , where the neutrino can be in any generation. Since this process violate B-L by two unit, it is naturally suppressed by the large mass scale of  $\chi$  and  $\eta$ . The decay rate can be given as:

$$\Gamma_{N_1} = |h_{1e}|^2 |f_{e\tau}|^2 \left( \frac{m^2}{M_\eta^2 M_\chi^2} \right)^2 \frac{M_{N_1}^5}{192\pi^3}, \quad (30)$$

where  $h_{1e}$  is the coupling of  $N_1$  to  $e^+$  and  $\eta^-$ , while  $f_{e\tau}$  is the coupling of  $\chi^-$  to  $e^-$  and  $\nu_\tau$ . Taking the B-L violating scale  $m = 10^9 \text{ GeV}$ ,  $M_\eta = 3 \times 10^{10} \text{ GeV}$  and  $M_\chi = 5 \times 10^{10} \text{ GeV}$ , as taken previously, the life time of  $N_1$  is found to be

$$\tau_{N_1} = 0.88 \times 10^{20} \text{ Sec} \left( \frac{O(1)}{|f_{e\tau}|^2} \right) \left( \frac{10^{-3}}{|h_{1e}|} \right)^2 \left( \frac{10 \text{ MeV}}{M_{N_1}} \right)^5. \quad (31)$$

Thus we see that  $\tau_{N_1} \gg \tau_0$ , where  $\tau_0 = 0.18 \times 10^{17} \text{ Sec}$ , is the age of the universe. Thus  $N_1$  can be a candidate of dark matter.

### VI. CONCLUSIONS

We proposed an extension of the Zee model with a conserved B-L global symmetry. The B-L symmetry is softly broken to give rise to the neutrino masses through one loop radiative correction. In contrast to the Zee model, in the present case the charged scalars flowing in the neutrino loop are super heavy. Therefore, these charged scalars could depart from thermal equilibrium in the early universe. As a result the CP violating decay of the super heavy charged particles, namely  $\eta^-$ , could generate a net baryon asymmetry through the leptogenesis route. Recall that the lepton asymmetry is generated without any B-L violation. This model then accommodate a sterile neutrino like dark matter  $N_1$  as its three body decay  $N_1 \rightarrow e^- e^+ \nu$  is suppressed by the large mass scale of  $\eta$  and  $\chi$ .

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